# Hard breakup of the deuteron into two $\Delta$ -isobars

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## Abstract

We study high energy photodisintegration of the deuteron into two $\Delta$-isobars at large center of mass angles within the QCD hard rescattering model (HRM). According to the HRM, the process develops in three main steps: the photon knocks a quark from one of the nucleons in the deuteron; the struck quark rescatters off a quark from the other nucleon sharing the high energy of the photon; then the energetic quarks recombine into two outgoing baryons which have large transverse momenta. Within the HRM, the cross section is expressed through the amplitude of $pn \to \Delta\Delta$ scattering which we evaluated based on the quark-interchange model of hard hadronic scattering. Calculations show that the angular distribution and the strength of the photodisintegration is mainly determined by the properties of the $pn \to \Delta\Delta$ scattering. We predict that the cross section of the deuteron breakup to $\Delta^{++}\Delta^{-}$ is 4-5 times larger than that of the breakup to the $\Delta^{+}\Delta^{0}$ channel. Also, the angular distributions for these two channels are markedly different. These can be compared with the predictions based on the assumption that two hard $\Delta$-isobars are the result of the disintegration of the preexisting $\Delta\Delta$ components of the deuteron wave function. In this case, one expects the angular distributions and cross sections of the breakup in both $\Delta^{++}\Delta^{-}$ and $\Delta^{+}\Delta^{0}$ channels to be similar.

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## I. INTRODUCTION

Hard nuclear processes provide an important testing ground for QCD degrees of freedom in nuclei. One of such processes is the high-energy large-angle photodisintegration of light nuclei. These reactions were intensively studied during the last two decades. The studies included the experiments on large center of mass (c.m.) angle break-up of the deuteron into the $pn$ pair[1–8] as well as break-up of the $^3$He nucleus into two high energy protons and a slow neutron[9]. The uniqueness of these reactions is in the effectiveness by which high momentum and energy are transferred to the NN system[10] at a given photon energy, $E_\gamma$. Namely at large and fixed values of the c.m. scattering angle, $s, -t \sim 2M_dE_\gamma$ which is by a factor of two larger than the invariant energy and transfered momenta achieved in hadronic interactions at the same incident energies.

The above mentioned reactions confirmed the prediction of quark-counting rule[11] according to which the energy dependence of the differential cross section at large c.m. scattering angles scales as $\frac{d\sigma}{dt} \sim s^{-11}$.

However, calculations of the absolute cross sections require a more detailed understanding of the dynamics of these processes. The considered theoretical models can be grouped by two distinctly different underlying assumptions made in the calculations[12]. The first assumes that the large c.m. angle nucleons are produced through the interaction of the incoming photon with a pre-existing hard two nucleon system in the nucleus[13–15]. The second approach is based on the assumption that the two high momentum nucleons are produced through a hard rescattering at the final state of the reaction[16–21].

In the hard rescattering model (HRM)[16] in particular, by explicitly introducing quark degrees of freedom, a parameter free cross section has been obtained for hard photodisintegration of the deuteron at $90^0$ c.m. angle [16, 17]. Also HRM's prediction of the hard breakup of two protons from the $^3$He nucleus[21] agreed reasonably well with the recent experimental data[9].

In the present work we extend the HRM approach to calculate hard breakup of the deuteron into two-$\Delta$-isobars produced at large angles in the $\gamma - d$ center of mass reference frame. In our estimates, we calculate the relative strength of $\gamma d \to \Delta^{++}\Delta^-$ and $\gamma d \to \Delta^+\Delta^0$ cross sections as they compare with the $\gamma d \to pn$ cross section.

The investigation of the production of two energetic $\Delta$-isobars from the deuteron has an important significance in probing possible non-nucleonic components in the deuteron wave function. Studies of non-nucleonic components of the deuteron have a rather long history. Already in the 1970's, the possible existence of the baryonic resonance components in the deuteron have been studied in potential and pion-exchange models (see e.g.[22–24, 24, 25]). They were also considered in quark-interchange[27] and chiral quark[26] models.

Among the all possible resonance components the $\Delta\Delta$ component has an interesting relation to the possible existence of the hidden color component in the deuteron wave function ( see e.g. Refs.[28–32]). This relation follows from the observation[28, 29] that in the regime in which the chiral symmetry is restored the color singled 6-quark configuration can be expressed through the superposition of $NN$, $\Delta\Delta$ and hidden-color components with relative normalizations fixed by the $SU(3)$ symmetry. Thus, experimental verification of the relative strength of the $NN$ to $\Delta\Delta$ component could shed light on the existence of hidden color components in the deuteron wave function. However, both components should be probed in hard nuclear processes in which case small inter-nucleon distances in the deuteron are probed. Our calculation in this case will allow us to asses the role of the hard rescattering in

these processes. It will allow us also to explore another venue for checking the basic mechanism of the high momentum transfer breakup of nuclei into two baryons. Our calculations result in the distinct predictions for angular distributions of the $\Delta$-isobar pair at large c.m. production angle as well as their relative strength compared with the production of the $pn$ pair at the same kinematics. Despite experimental challenges associated with the investigation of two $\Delta$-isobar breakup of the deuteron[34], there are ongoing efforts in performing such experiment at Jefferson Lab[35, 36] which we hope will allow to verify our predictions.

## II. HARD RESCATTERING MODEL

We consider the photoproduction of two baryons, $B_1$ and $B_2$, in the reaction,

$$\gamma + d \rightarrow B_1 + B_2 \tag{1}$$

in which the baryons are produced at large angles in the $\gamma - d$ center of mass reference frame.

According to the HRM, the large angle breakup of the NN system proceeds through the knock-out of a valence quark from one of the nucleons with subsequent hard rescattering of the struck-quark with a valence quark of the second nucleon. The two quarks then recombine with the spectator systems of nucleons forming two emerging baryons with large transverse momenta. The hard rescattering provides the mechanism of sharing the photon's energy among two final baryons.

The invariant amplitude of the photodisintegration Eq.(1) is calculated by applying Feynman diagram rules to diagrams similar to Fig.1. During the calculation we introduce undetermined quark wave functions of baryons to account for the transition of the initial nucleons to the quark-spectator systems, and also for the recombination of the final state quarks with these spectator systems into the final two baryon system.

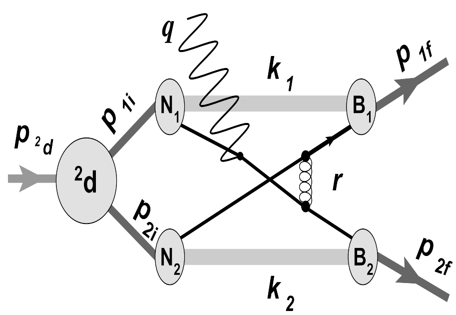


FIG. 1: Deuteron photodisintegration according to the HRM

Fig.1 displays the chosen independent momenta for three loop integration involved in the invariant amplitude. Two major approximations simplify further calculations. First, using the fact that the struck quark is very energetic we treat it on its mass shell. Then the struck quark's propagator is evaluated at it's pole value at such magnitudes of nucleon momenta that maximize the deuteron wave function. These approximations allow us to factorize the invariant amplitude into three distinguished parts. The first, representing the transition amplitude of the deuteron into the $(pn)$ system, which can be evaluated using a realistic deuteron wave function. The second is the amplitude of photon-quark interaction, and the

third term represents the hard rescattering of the struck quark with recombination into a two large transverse momentum baryonic system. Combined with the initial state nucleon wave functions, the rescattering part is expressed through the quark-interchange (QI) amplitude of $pn \to B_1B_2$ scattering. Details of the derivation are given in Refs.[16, 21]. After the above mentioned factorization is made, the overall invariant amplitude of $\gamma d \to B_1B_2$ reaction can be expressed as follows:

$$\langle \lambda_{1f}, \lambda_{2f} \mid \mathcal{M} \mid \lambda_\gamma, \lambda_d \rangle = ie[\lambda_\gamma] \times$$
$$\Bigg\{ \sum_{i \in N_1} \sum_{\lambda_{2i}} \int \frac{Q_i^{N_1}}{\sqrt{2s'}} \langle \lambda_{2f}; \lambda_{1f} \mid T^{QI}_{(pn \to B_1B_2),i}(s, t_N) \mid \lambda_\gamma; \lambda_{2i} \rangle \Psi_{\rm d}^{\lambda_d}(p_{1i}, \lambda_\gamma; p_{2i}, \lambda_{2i}) \frac{d^2p_\perp}{(2\pi)^2}$$
$$+ \sum_{i \in N_2} \sum_{\lambda_{1i}} \int \frac{Q_i^{N_2}}{\sqrt{2s'}} \langle \lambda_{2f}; \lambda_{1f} \mid T^{QI}_{(pn \to B_1B_2),i}(s, t_N) \mid \lambda_{1i}; \lambda_\gamma \rangle \Psi_{\rm d}^{\lambda_d}(p_{1i}, \lambda_{1i}; p_{2i}, \lambda_\gamma) \frac{d^2p_\perp}{(2\pi)^2} \Bigg\} \quad (2)$$

where $\lambda_\gamma$,$\lambda_d$, $\lambda_{1f}$ and $\lambda_{2f}$ are the helicities of the photon, deuteron and the two outgoing baryons respectively. Here $\Psi_{\rm d}^{\lambda_d}(p_{1i}, \lambda_{1i}; p_{2i}, \lambda_{2i})$ is the $\lambda_d$-helicity light-cone deuteron wave function defined in the $q_+ = 0$ reference frame. The initial light-cone momenta of the nucleons in the deuteron are $p_{1i} = (\alpha_{1i} = \frac{1}{2}, p_{1i\perp} = -p_\perp)$ and $p_{2i} = (\alpha_{2i} = \frac{1}{2}, p_{2i\perp} = p_\perp)$ with $\lambda_{1i}$ and $\lambda_{2i}$ being their helicities respectively. The $\frac{1}{\sqrt{s'}}$ factor with $s' = s - M_d^2$ comes from the energetic propagator of the struck quark before its rescattering. The squares of the total invariant energy as well as the momentum transfer are defined as follows:

$$\begin{aligned} s &= (q + p_d)^2 = (p_{1f} + p_{2f})^2 = 2E_\gamma^{lab} M_d + M_d^2 \\ t &= (p_{1f} - q)^2 = (p_{2f} - p_d)^2 \end{aligned} \quad (3)$$

where $q$, $p_d$, $p_{1f}$ and $p_{2f}$ are the four-momenta of the photon, deuteron and two outgoing baryons respectively. The lab energy of the photon is defined by $E_\gamma^{lab}$, and $M_d$ is the mass of the deuteron. The transfer momentum, $t_N$ in the rescattering amplitude in Eq.(2) is defined as:

$$t_N = (p_{1f} - p_{1i} - q)^2 = (p_{2f} - p_{2i}) \approx (p_{2f} - \frac{p_d}{2})^2 = \frac{t}{2} + \frac{m_{B2}^2}{2} - \frac{M_d^2}{4}, \quad (4)$$

where the approximation in the right hand side follows from the assumption that the magnitudes of light-cone momentum fractions of bound nucleons dominating in the scattering amplitude are $\alpha_{1i} = \alpha_{2i} = \frac{1}{2}$, and that the transverse momenta of these nucleons are negligible as compared to the momentum transfer in the reaction, $p_\perp^2 \ll |t_N|, |u_N|$.

In Eq.(2) the following expression

$$Q_i \langle \lambda_{2f}; \lambda_{1f} \mid T^{QI}_{(pn \to B_1B2),i}(s, t_N) \mid \lambda_{1i}; \lambda_{2i} \rangle \quad (5)$$

represents the quark-charge weighted QI amplitude of $pn \to B_1B_2$ hard exclusive scattering. The factor $Q_i$ corresponds to the charge (in $e$ units) of the quark that interacts with the incoming photon. In a further approximation we factorize the hard rescattering amplitude from the integral since the momentum transfer entering in $T_{(pn \to B_1B_2),i}(s, t_N)$ significantly exceeds the Fermi momentum of the nucleon in the deuteron. Also, after calculating the overall quark-charge factors, the QI scattering amplitudes are identified with the $NN \to B_1B_2$ helicity amplitudes as follows:

$$\langle \lambda_{2f}; \lambda_{1f} \mid T^{QI}_{pn \to B_1B_2}(s, t_N) \mid \lambda_{1i}; \lambda_{2i} \rangle = \phi_j(s, \theta^N_{c.m.}), \quad (6)$$

where $\theta^N_{c.m.}$ is the effective center of mass angle defined for given $s$ and $t_N$.

The differential cross section for unpolarized scattering is obtained through:

$$\frac{d\sigma_{\gamma d\to B_1 B_2}}{dt} = \frac{1}{16\pi}\frac{1}{(s-M_d^2)}|\bar{\mathcal{M}}|^2_{\gamma d\to B_1B_2} \tag{7}$$

where

$$|\bar{\mathcal{M}}|^2_{\gamma d\to B_1B_2} = \frac{1}{3}\frac{1}{2}\sum_{\lambda_{1f},\lambda_{2f},\lambda_\gamma,\lambda_d} |\langle\lambda_{1f},\lambda_{2f}\mid\mathcal{M}\mid\lambda_\gamma,\lambda_d\rangle|^2, \tag{8}$$

with the invariant amplitude square averaged by the number of helicity states of the deuteron and photon.

## III. CROSS SECTION OF THE $\gamma + d \to pn$ BREAKUP REACTION

We derive the amplitude of the breakup of the deuteron into the $pn$ pair from Eq.(2) by introducing the independent helicity amplitudes of $pn$ elastic scattering Eq.(A2) and by separating the quark-charge factors into $\hat{Q}^{N_1}$ and $\hat{Q}^{N_2}$ which correspond to the scattering of the photon off the quark of the first and the second nucleons in the deuteron. Then, for Eq.(8) one obtains:

$$\begin{aligned}|\bar{\mathcal{M}}|^2 \;=\; & \frac{1}{2}\frac{1}{3}\frac{e^2}{2s'}\Big[S_{12}\Big\{|(\hat{Q}^{N_1}+\hat{Q}^{N_2})\phi_1|^2+|(\hat{Q}^{N_1}+\hat{Q}^{N_2})\phi_2|^2\Big\}\\ & + S_{34}\Big\{|\hat{Q}^{N_1}\phi_3+\hat{Q}^{N_2}\phi_4|^2+|\hat{Q}^{N_1}\phi_4+\hat{Q}^{N_2}\phi_3|^2\Big\}\\ & + 2S_0|(\hat{Q}^{N_1}+\hat{Q}^{N_2})\phi_5|^2\Big],\end{aligned} \tag{9}$$

where the light-cone spectral functions of the deuteron are defined as follows:

$$\begin{aligned} S_{12} &= \sum_{\lambda=-1}^{1}\sum_{(\lambda_1=\lambda_2=-\frac{1}{2})}^{\frac{1}{2}}\left|\int\Psi_{\rm d}^{\lambda_d}(p_1,\lambda_1;p_2,\lambda_2)\frac{d^2p_\perp}{(2\pi)^2}\right|^2,\\ S_{34} &= \sum_{\lambda=-1}^{1}\sum_{(\lambda_1=-\lambda_2=-\frac{1}{2})}^{\frac{1}{2}}\left|\int\Psi_{\rm d}^{\lambda_d}(p_1,\lambda_1;p_2,\lambda_2)\frac{d^2p_\perp}{(2\pi)^2}\right|^2,\\ S_0 &= S_{12}+S_{34}.\end{aligned} \tag{10}$$

Eq.(9) can be further simplified if we assume (see e.g.[37]) that $\phi_3 \approx \phi_4$, as well as $S_{12} \approx S_{34} = \frac{S_0}{2}$, which results in:

$$|\bar{\mathcal{M}}|^2 \;=\; \frac{1}{2}\frac{1}{3}\frac{e^2}{2s'}Q^2_{F,pn}\frac{S_0}{2}\Big[|\phi_1|^2+|\phi_2|^2+|\phi_3|^2+|\phi_4|^2+|\phi_4|^2+4|\phi_5|^2\Big]. \tag{11}$$

Using the expression of the differential cross section of elastic $pn$ scattering:

$$\frac{d\sigma^{NN\to NN}(s,\theta^N_{c.m.})}{dt} = \frac{1}{16\pi}\frac{1}{s(s-4m_N^2)}\frac{1}{2}(|\phi_1|^2+|\phi_2|^2+|\phi_3|^2+|\phi_4|^2+4|\phi_5|^2), \tag{12}$$

and the relation between the light-cone and non-relativistic deuteron wave functions[16, 38–40] at small internal momenta: $\Psi_d(\alpha, p_\perp) = (2\pi)^{\frac{3}{2}}\Psi_{d,NR}(p)\sqrt{m_N}$ in Eq.(9), for the differential cross section on obtains from Eq.(7):

$$\frac{d\sigma^{\gamma d\to pn}(s,\theta_{c.m.})}{dt} = \frac{\alpha Q^2_{F,pn} 8\pi^4}{s'}\frac{d\sigma^{pn\to pn}(s,\theta^N_{c.m.})}{dt}\bar{S}_{0,NR}, \tag{13}$$

where we neglected the difference between $4m_N^2$ and $M_d^2$. Here the averaged non relativistic spectral function of the deuteron is defined as follows:

$$\bar{S}_{0,NR} = \frac{1}{3}\sum_{\lambda=-1}^{\lambda=1}\sum_{\lambda_1,\lambda_2=-\frac{1}{2}}^{\frac{1}{2}}\left|\int \Psi^{\lambda_d}_{\rm d,NR}(\alpha=\frac{1}{2}, p_\perp, \lambda_1; \alpha=\frac{1}{2}, -p_\perp, \lambda_2)\sqrt{m_N}\frac{d^2p_\perp}{(2\pi)^2}\right|^2, \tag{14}$$

where $\Psi_{d,NR}$ is the non relativistic deuteron wave function, which can be calculated using realistic $NN$ interaction potentials.

The quark-charge factor, $Q_{F,pn} = \frac{1}{3}$[16] accounts for the amount of the effective charge exchanged between the proton and the neutron in the rescattering. It is estimated by counting all the possible quark-exchanges within the $pn$ pair weighted with the charge of one of the exchanged quarks (for more details see Appendix B). The result in Eq.(13) is remarkably simple and contains no free parameters. It can be evaluated using the experimental values of the differential cross section of the elastic $pn$ scattering, $\frac{d\sigma^{pn\to pn}(s,\theta^N_{c.m.})}{dt}$. The angle $\theta^N_{c.m.}$ entering in the $pn\to pn$ cross section is the center of mass angle of the scattering corresponding to the $NN$ elastic reaction at $s$ and $t_N$. It is related to $\theta_{c.m.}$ of the $pn$ photodisintegration by [21]:

$$cos(\theta^N_{c.m.}) = 1 - \frac{(s-M_d^2)}{2(s-4m_N^2)}\frac{(\sqrt{s}-\sqrt{s-4m_N^2}cos(\theta_{c.m.}))}{\sqrt{s}} + \frac{4m_N^2-M_d^2}{2(s-4m_N^2)}. \tag{15}$$

It is worth mentioning that as it follows from the equation above, $\theta_{c.m.} = 90^0$ photodisintegration will correspond to the $\theta^N_{c.m.} = 60^0$ hard $pn$ elastic rescattering at the final state of the reaction.

## IV. CROSS SECTION OF THE $\gamma d \to \Delta\Delta$ BREAKUP REACTION

We use an approach similar to that in Sec.III to derive the invariant amplitude of the $\gamma d\to\Delta\Delta$ reactions. In this case Eq.(2) requires an input of the helicity amplitudes of the corresponding $pn\to\Delta\Delta$ scattering. One has a total 32 independent helicity amplitudes for this scattering. To simplify further our derivations, we will restrict ourselves by considering only the seven helicity conserving amplitudes given in Eq.(A3). Using these amplitudes in Eq.(2) and separating the quark-charge factors into $\hat{Q}^{N_1}$ and $\hat{Q}^{N_2}$, similar to Eq.(9) one obtains

$$\begin{aligned}|\bar{\mathcal{M}}|^2{}_{\gamma d\to\Delta\Delta} = {} & \frac{1}{2}\frac{1}{3}\frac{e^2}{2s'}\Big[S_{12}\Big\{|(\hat{Q}^{N_1}+\hat{Q}^{N_2})\phi_1|^2 + |(\hat{Q}^{N_1}+\hat{Q}^{N_2})\phi_6|^2 + |(\hat{Q}^{N_1}+\hat{Q}^{N_2})\phi_7|^2\Big\} \\ & + S_{34}\Big\{|\hat{Q}^{N_1}\phi_3+\hat{Q}^{N_2}\phi_4|^2 + |\hat{Q}^{N_1}\phi_4+\hat{Q}^{N_2}\phi_3|^2 \\ & + |\hat{Q}^{N_1}\phi_8+\hat{Q}^{N_2}\phi_9|^2 + |\hat{Q}^{N_1}\phi_9+\hat{Q}^{N_2}\phi_8|^2\Big\}\Big],\end{aligned} \tag{16}$$

where $S_{12}$ and $S_{34}$ are defined in Eq.(10). Similar to the previous section, we simplify further the above expression assuming that all helicity conserving amplitudes are of the same order of magnitude. Assuming also that $S_{12} \approx S_{34} \approx \frac{S_0}{2}$, we obtain

$$|\bar{\mathcal{M}}|^2 = \frac{1}{2}\frac{1}{3}\frac{e^2}{2s'}Q_{F,\Delta\Delta}\frac{S_0}{2}\left[|\phi_1|^2 + |\phi_3|^2 + |\phi_4|^2 + |\phi_4|^2 + |\phi_6|^2 + |\phi_7|^2 + |\phi_8|^2 + |\phi_9|^2\right], \quad (17)$$

where $Q_{F,\Delta\Delta} = \hat{Q}^{N_1} + \hat{Q}^{N_2} = \frac{1}{3}$ is obtained by using the same approach as for the case of the $pn$ breakup in Sec.III. Using now the expression of the differential cross section of $pn \to \Delta\Delta$ scattering,

$$\frac{d\sigma^{pn\to\Delta\Delta}(s, \theta^N_{c.m.})}{dt} = \frac{1}{16\pi}\frac{1}{(s - 4m_N^2)}\frac{1}{2}\left[|\phi_1|^2 + |\phi_3|^2 + |\phi_4|^2 + |\phi_4|^2 + |\phi_6|^2 + |\phi_7|^2 + |\phi_8|^2 + |\phi_9|^2\right] \quad (18)$$

as well as the relation between light-cone and non relativistic deuteron wave function discussed in Sec.III, from Eq.(7) we obtain the following expression for the differential cross section of the $\gamma d \to \Delta\Delta$ scattering:

$$\frac{d\sigma^{\gamma d\to\Delta\Delta}(s, \theta_{c.m.})}{dt} = \frac{\alpha Q^2_{F,\Delta\Delta} 8\pi^4}{s'}\frac{d\sigma^{pn\to\Delta\Delta}(s, \theta^N_{c.m.})}{dt}\bar{S}_{0,NR}, \quad (19)$$

where $\bar{S}_{0,NR}$ is given in Eq.(14). The effective c.m. angle $\theta^N_{c.m.}$ entering in the argument of the differential cross section of $pn \to \Delta\Delta$ reaction can be calculated by using Eqs. (3) and (4) to obtain

$$cos\theta^N_{c.m.} = \frac{1}{\sqrt{(s - 4m_N^2)(s - 4m_\Delta^2)}}\left[s - \frac{M_d^2 - 4m_N^2}{2} - \frac{s - M_d^2}{2\sqrt{s}}\left(\sqrt{s} - \sqrt{s - 4m_\Delta^2}cos\theta_{c.m.}\right)\right]. \quad (20)$$

As it follows from Eq.(19), provided there are enough experimental data on high momentum transfer $pn \to \Delta\Delta$ differential cross sections, the $\gamma d \to \Delta\Delta$ cross section can be computed without introducing an adjustable free parameter. However, there are no experimental data on hard exclusive $pn \to \Delta\Delta$ reactions with sufficient accuracy that would allow us to make quantitative estimates based on Eq.(19). Instead, in the next section we will attempt to make quantitative predictions based on the quark-interchange framework of hard scattering.

## V. ESTIMATES OF THE RELATIVE STRENGTH OF THE $\Delta\Delta$ BREAKUP REACTIONS.

Our further calculations are based on the experimental observation[41] that the quark-interchange[42] represents the dominant mechanism of hard exclusive scattering of baryons that carry valence quarks with common flavor. The quark-interchange mechanism however will not allow us to calculate the absolute cross sections. Instead, we expect that its predictions will be more reliable for the ratios of the differential cross sections for different exclusive channels.

As an illustration of the reliability of calculations of cross section ratios in the QI model, in Fig.2 we compare the QI predictions for the ratios of $pn$ to $pp$ differential cross sections at $90^0$ c.m. scattering. Here, we compare predictions based on SU(6)[43, 44] and diquark[45] symmetry assumptions for the valence quark wave function of the nucleons. As comparison shows one achieves a rather reasonable agreement with the data without any additional normalization parameter. Based on this, we now estimate the ratio of the differential cross sections of $\gamma d \to \Delta\Delta$ to the $\gamma d \to pn$ cross sections. We use both SU(6) and diquark-symmetry quark wave functions of the nucleon and $\Delta$-isobars (see Appendix B) in the calculation of the $pn \to \Delta\Delta$ amplitudes.

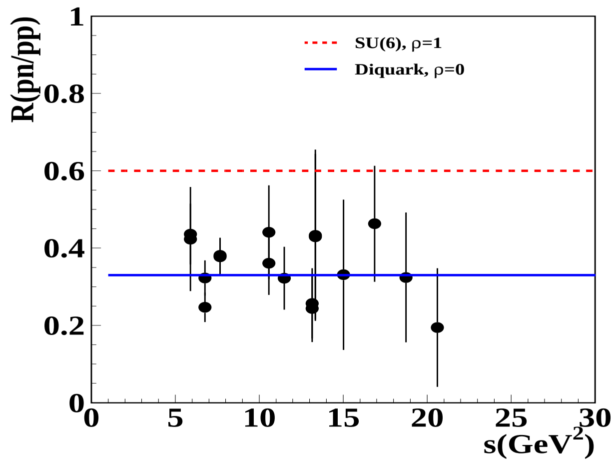


FIG. 2: (Color online) Ratio of the $pn \to pn$ to $pp \to pp$ elastic differential cross sections as a function of $s$ at $\theta^N_{c.m.} = 90^0$.

To calculate the photodisintegration amplitudes we go back to Eqs.(9) and (16) and evaluate the quark-charge factors using SU(6) or diquark symmetries of the valence quark wave functions of baryons. For this we separate the $t$ and $u$ channels in the helicity amplitudes:

$$\phi_i(s, \theta^N_{c.m.}) = \phi^t_i(s, \theta^N_{c.m.}) + \phi^u_i(s, \theta^N_{c.m.}) \tag{21}$$

and then treat the charge factors for the given nucleon $N$ as:

$$\hat{Q}^N \phi_l = Q^{t,N}_i \phi^t_l + Q^{u,N}_i \phi^u. \tag{22}$$

This yields the following expression for the photodisintegration amplitude of Eq.(2) :

$$\begin{aligned} \langle \lambda_{1f}, \lambda_{2f} \mid \mathcal{M} \mid \lambda_\gamma, \lambda_d \rangle \;&=\; ie[\lambda_\gamma] \times \Bigg\{ \sum_{\lambda_{2i}} \frac{1}{\sqrt{2s'}} \left[ Q^{tN_1}_i \phi^t_i + Q^{uN_1}_i \phi^u_i \right]_{\lambda_{2i}} \int \Psi^{\lambda_d}_{\rm d}(p_1, \lambda_\gamma; p_2, \lambda_{2i}) \frac{d^2 p_\perp}{(2\pi)^2} \\ &+ \sum_{\lambda_{1i}} \frac{1}{\sqrt{2s'}} \left[ Q^{tN_2}_i \phi^t_i + Q^{uN_2}_i \phi^u_i \right]_{\lambda_{1i}} \int \Psi^{\lambda_d}_{\rm d}(p_1, \lambda_{1i}; p_2, \lambda_\gamma) \frac{d^2 p_\perp}{(2\pi)^2} \Bigg\}. \end{aligned} \tag{23}$$

### A. $\gamma d \to pn$ scattering

For the $\gamma d \to pn$ amplitude, the charge factors calculated for the helicity conserving amplitudes according to the QI framework yield for both SU(6) and diquark models (see

Appendix B)

$$
\begin{aligned}
Q_j^{tN_1} &= Q_j^{tN_2} = \frac{Q_{F,pn}}{2} \\
Q_j^{uN_1} &= -2Q_j^{uN_2} = 2Q_{F,pn}
\end{aligned} \tag{24}
$$

with $Q_{F,pn} = \frac{1}{3}$ and independent of j. Using these relations in Eq.(22), from Eqs.(23) and (9) one obtains

$$
|\bar{\mathcal{M}}|^2_{\gamma d\longrightarrow pn} = \frac{e^2}{6\cdot 2s'}Q^2_{F,pn}\left\{S_{12}\phi_1^2 + S_{34}\left[\left(\frac{\phi_3^t+\phi_4^t}{2}+2\phi_4^u-\phi_3^u\right)^2+\left(\frac{\phi_4^t+\phi_3^t}{2}+2\phi_3^u-\phi_4^u\right)^2\right]\right\}, \tag{25}
$$

where the different predictions of SU(6) and diquark models follow from the different predictions for the $pn\to pn$ helicity conserving amplitudes given in Eq.(B8).

### B. $\gamma d\to\Delta^+\Delta^0$ scattering

The calculation for the $\gamma d\to\Delta^+\Delta^0$ amplitude yields the same quark-charge factors as for the $\gamma d\to pn$ reactions in Eq.(24). Using the helicity amplitudes of the $pn\to\Delta^+\Delta^0$ scattering from Eq.(B9) and the expressions for the photodisintegration amplitudes from Eqs.(23,16) one obtains

$$
\begin{aligned}
|\bar{\mathcal{M}}|^2_{\gamma d\longrightarrow\Delta^+\Delta^-} = &\frac{1}{6}\frac{e^2}{2s'}Q^2_{F,\Delta\Delta}\left\{S_{12}\left[|\phi_1|^2+|\phi_6|^2+|\phi_7|^2\right]\right. \\
&+S_{34}\left[\left(\frac{\phi_3^t+\phi_4^t}{2}+2\phi_4^u-\phi_3^u\right)^2+\left(\frac{\phi_4^t+\phi_3^t}{2}+2\phi_3^u-\phi_4^u\right)^2\right. \\
&\left.\left.+\left(\frac{\phi_8^t+\phi_9^t}{2}+2\phi_9^u-\phi_8^u\right)^2+\left(\frac{\phi_9^t+\phi_8^t}{2}+2\phi_8^u-\phi_9^u\right)^2\right]\right\},
\end{aligned} \tag{26}
$$

where the different predictions of SU(6) and diquark models follow from the different predictions for the $pn\to\Delta^+\Delta^0$ helicity conserving amplitudes given in Eq.(B9).

### C. $\gamma d\to\Delta^{++}\Delta^-$ scattering

For the charge factors in the $\gamma d\to\Delta^{++}\Delta^-$ scattering within the quark-interchange approximation from Appendix B we obtain:

$$
-Q^{tN_1} = \frac{Q^{tN_2}}{2} = Q_{F,\Delta\Delta} = \frac{1}{3}. \tag{27}
$$

Inserting these charge factors in Eqs.(23,16) one obtains for the photodisintegration amplitude:

$$
\begin{aligned}
|\bar{\mathcal{M}}|^2_{\gamma d\longrightarrow\Delta^{++}\Delta^-} = &\frac{1}{6}\frac{e^2}{2s'}Q^2_{F,\Delta\Delta}\left\{S_{12}\left(|\phi_1|^2+|\phi_6|^2+|\phi_7|^2\right)\right. \\
&\left.+S_{34}\left[(2\phi_3-\phi_4)^2+(2\phi_4-\phi_3)^2+5|\phi_8|^2\right]\right\}.
\end{aligned} \tag{28}
$$

where predictions for the helicity conserving amplitudes of $pn\to\Delta^{++}\Delta^-$ are given in Eq.(B10).

### D. Numerical Estimates

Using Eqs.(25), (26) and (28) with the baryonic helicity amplitudes calculated in Appendix B we estimate the ratio $R(\theta_{c.m.})$ of the $\gamma d \to \Delta\Delta$ to $\gamma d \to pn$ differential cross sections at given $s$ and $\theta_{c.m.}$ angle. For simplicity we consider the kinematics in which $s >> 4m_\Delta^2$, which allows to approximate both Eqs.(15) and (20) to,

$$cos\theta^N_{c.m} \approx \frac{1 + cos\theta_{c.m.}}{2}. \tag{29}$$

Before considering any specific model for angular distribution, one can make two general statements about the properties of the photodisintegration amplitude. First, that from the absence of the $u$ channel scattering in the $pn \to \Delta^{++}\Delta^-$ helicity amplitudes (see Eq.(B10)), one observes that $R(\theta_{c.m.})$ can not be a uniform function of $\theta_{c.m}$. Second, that independent of the choice of SU(6) or diquark models, the $\gamma d \to \Delta^{++}\Delta^-$ cross section is always larger than the cross section of the $\gamma d \to \Delta^+\Delta^-$ reaction.

We quantify the above observations by parameterizing the angular function $f(\theta^N_{c.m.})$, which enters in Eqs.(B8,B9,B10), in the following form[37, 45]:

$$f(\theta) = \frac{1}{sin(\theta)^2(1 - cos(\theta))^2} \tag{30}$$

known to describe reasonably well the elastic $pp$ and $pn$ scattering cross sections.

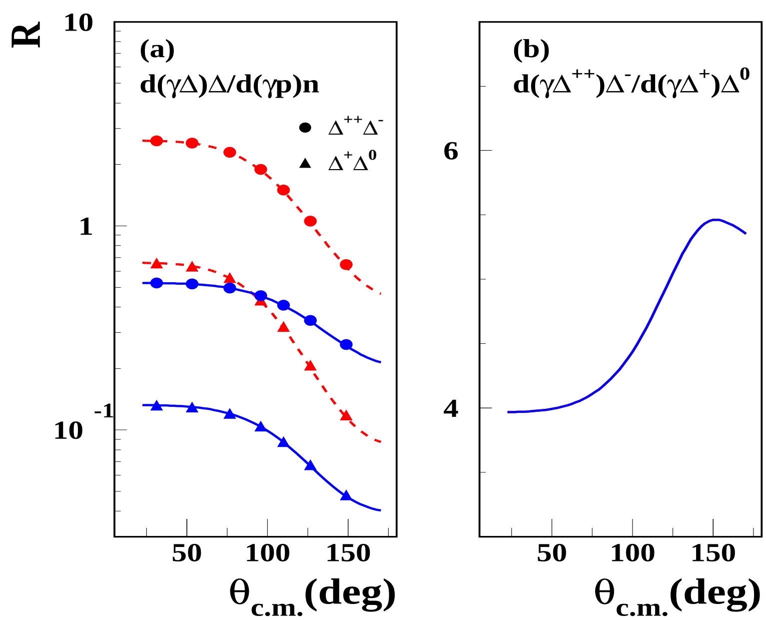


FIG. 3: (Color online) (a)Ratio of the $\gamma d \to \Delta\Delta$ to $\gamma d \to pn$ differential cross sections and (b) ratio of the $\gamma d \to \Delta^{++}\Delta^-$ to $\gamma d \to \Delta^+\Delta^0$ differential cross sections as a function of $\theta_{c.m.}$.

Magnitudes of the ratio $R$ at $\theta_{c.m} = 90^0$ are given in Table I, while the angular dependencies (solid curves for diquark model and dashed curves for SU(6) model) are presented in Fig.3(a). They clearly show strong angular anisotropy and the excess (by a factor of 4-5) of the $\Delta^{++}\Delta^-$ breakup cross section relative to the cross section of the $\Delta^+\Delta^0$ breakup (Fig.3(b)). Our calculations show that the ratio of the $\gamma d \to \Delta\Delta$ to $\gamma d \to pn$ cross sections

is very sensitive to the choice of SU(6) or diquark models of the wave functions. However, because of the absence of isosinglet two-quark state in the $\Delta$ wave functions, the $\rho$ parameter dependence that characterizes the choice of SU(6) or diquark models in the baryons wave functions is factorized and enters only in the normalization factor of the $pn \to \Delta\Delta$ helicity amplitudes. As a result, the ratio of the $\gamma d \to \Delta^{++}\Delta^{-}$ to $\gamma d \to \Delta^{+}\Delta^{0}$ cross sections (Fig.3b) is independent of the choice between SU(6) and diquark models for the baryons wave functions.

Finally, it is worth discussing how our calculations compare with the predictions of models in which the production of two $\Delta$'s is a result of the breakup of the pre-existing $\Delta\Delta$ component of the deuteron wave function. In this case, the final state interaction is dominated by soft scattering of two $\Delta$'s in the final state which will induce similar angular distributions for both $\Delta^{++}\Delta^{-}$ and $\Delta^{+}\Delta^{0}$ channels (see e.g.,[48, 49]). As a result, we expect essentially the same angular distribution for both $\Delta^{++}\Delta^{-}$ and $\Delta^{+}\Delta^{0}$ production channels. Also, because of the deuteron being an isosinglet, the probabilities of finding preexisting $\Delta^{++}\Delta^{-}$ and $\Delta^{+}\Delta^{0}$ are equal. For coherent hard breakup of the preexisting $\Delta$'s we will obtain the same cross section for both the $\Delta^{++}\Delta^{-}$ and the $\Delta^{+}\Delta^{0}$ channels.

| | $R(90^o)$ | |
|---|---|---|
| $\gamma d \to BB$ | SU(6) | Diquark |
| $\gamma d \to \Delta^{+}\Delta^{0}$ | 0.47 | 0.11 |
| $\gamma d \to \Delta^{++}\Delta^{-}$ | 2.01 | 0.47 |

TABLE I: Strength of $\Delta\Delta$ channels relative to $pn$ in deuteron photodisintegration at $\theta_{c.m} = 90^o$.

One interesting scenario for probing the preexisting $\Delta$'s in the deuteron is using the decomposition of the deuteron wave function, in the chiral symmetry restored limit, into the nucleonic and non-nucleonic components in the following form[28–30]:

$$\Psi_{T=0,S=1} = (\frac{1}{9})^{\frac{1}{2}}\Psi_{NN} + (\frac{4}{45})^{\frac{1}{2}}\Psi_{\Delta\Delta} + (\frac{4}{5})^{\frac{1}{2}}\Psi_{CC}, \tag{31}$$

where $\Psi_{CC}$ represents the hidden color component of $T = 0$ and $S = 1$ six-quark configuration. Since $\Delta^{++}\Delta^{-}$ and $\Delta^{+}\Delta^{0}$ components enter with equal probability in the total isospin $T = 0$ configuration, one expects close ($\approx 0.8$) strengths for deuteron breakup to $\Delta^{++}\Delta^{-}$ or $\Delta^{+}\Delta^{0}$ channels as compared to the strength of the deuteron breakup into the $pn$ pair. This result should be compared with the similar ratios presented in Table I from HRM and with the HRM angular distributions in Fig. 3.

It is worth noting that HRM can be applied for calculation of the large angle photoproduction of any given two baryonic resonances. In all cases the model will be sensitive to the valence quark wave function of the baryons as well as to the effective color charge factors entering in the scattering amplitude. One such possibility is the large center of mass angle photoproduction of the $NN^*$ pair which will allow us to evaluate the role of the rescattering in reactions aimed at probing the $NN*$ component of the deuteron wave function. Note that such a process will not interfere with the amplitude of $\Delta\Delta$ production at large center of mass angles, since the decay products of the produced resonances occupy distinctly different phase spaces in the final state of the reaction.

## VI. SUMMARY

We extended the hard rescattering model of large c.m. angle photodisintegration of a two-nucleon system to account for the production of two $\Delta$-isobars. The HRM allows to express the cross section of $\gamma d \to pn$ and $\gamma d \to \Delta\Delta$ reactions through the large c.m. angle differential cross section of $pn \to pn$ and $pn \to \Delta\Delta$ scattering amplitudes.

Because of lack of experimental information on $pn \to \Delta\Delta$ scattering, we further applied the quark-interchange model to calculate the strength of the $\gamma d \to \Delta\Delta$ cross section relative to the cross section of $\gamma d \to pn$ breakup reaction. We predicted a significantly larger strength for the $\Delta^{++}\Delta^{-}$ channel of breakup as compared to the $\Delta^{+}\Delta^{0}$ channel which is related to the relative strength of the $pn \to \Delta^{++}\Delta^{-}$ and $pn \to \Delta^{+}\Delta^{0}$ scatterings. Because of the different angular dependences of these hadronic amplitudes, we also predicted a significant difference between the angular dependences of photoproduction cross sections in $\Delta^{++}\Delta^{-}$ and $\Delta^{+}\Delta^{0}$ channels.

These results can be compared with the prediction of the models in which two $\Delta$'s are produced due to the coherent breakup of the $\Delta\Delta$ component of the deuteron wave function. In this case one expects essentially similar angular distributions and strengths for the $\Delta^{++}\Delta^{-}$ and $\Delta^{+}\Delta^{0}$ breakup channels.

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### Appendix A: Baryon-Baryon Scattering Helicity Amplitudes

We are using helicity states to label the entries of the photodisintegration and the baryon-baryon scattering matrices. The number of independent helicity amplitudes for a given $ab \to cd$ processes can be expressed through the total spin of the scattering particles as follows[46, 47]:

$$N = \frac{1}{2} \cdot (2s_a + 1)(2s_b + 1)(2s_c + 1)(2s_d + 1) \tag{A1}$$

where $s_i$ is the total spin of particle i and for the photon we replace $(s_i + 1)$ by 2. The factor $\frac{1}{2}$ follows from the constraint due to the parity conservation. For elastic scattering, there is a further reduction in $N$ due to time reversal invariance, and if the scattering particles are identical, or lie in the same isospin multiplet, the number of independent helicity amplitudes is reduced further[46, 47]. For the $pn$ elastic scattering case, out of the possible 16 helicity amplitudes only five are independent[47] for which we use the following notations:

$$\begin{aligned}
\left\langle +\frac{1}{2}, +\frac{1}{2} \,|T|\, +\frac{1}{2}, +\frac{1}{2} \right\rangle &= \phi_1 \\
\left\langle +\frac{1}{2}, -\frac{1}{2} \,|T|\, +\frac{1}{2}, -\frac{1}{2} \right\rangle &= \phi_3
\end{aligned}$$

$$
\begin{aligned}
\left\langle -\frac{1}{2}, +\frac{1}{2} \left| T \right| +\frac{1}{2}, -\frac{1}{2} \right\rangle &= \phi_4 \\
\left\langle -\frac{1}{2}, -\frac{1}{2} \left| T \right| +\frac{1}{2}, +\frac{1}{2} \right\rangle &= \phi_2 \\
\left\langle -\frac{1}{2}, +\frac{1}{2} \left| T \right| +\frac{1}{2}, +\frac{1}{2} \right\rangle &= \phi_5,
\end{aligned}
\tag{A2}
$$

For the $pn \to \Delta\Delta$ scattering amplitude, we have from Eq.(A1), $N$=(2)(2)(4)(4)/2=32 independent helicity amplitudes. We use the following notations for the helicity conserving independent amplitudes of $pn \to \Delta\Delta$ scattering:

$$
\begin{aligned}
\left\langle +\frac{1}{2}, +\frac{1}{2} \left| T \right| +\frac{1}{2}, +\frac{1}{2} \right\rangle &= \phi_1 \\
\left\langle +\frac{1}{2}, -\frac{1}{2} \left| T \right| +\frac{1}{2}, -\frac{1}{2} \right\rangle &= \phi_3 \\
\left\langle -\frac{1}{2}, +\frac{1}{2} \left| T \right| +\frac{1}{2}, -\frac{1}{2} \right\rangle &= \phi_4 \\
\left\langle +\frac{3}{2}, -\frac{1}{2} \left| T \right| +\frac{1}{2}, +\frac{1}{2} \right\rangle &= \phi_6 \\
\left\langle -\frac{1}{2}, +\frac{3}{2} \left| T \right| +\frac{1}{2}, +\frac{1}{2} \right\rangle &= \phi_7 \\
\left\langle +\frac{3}{2}, -\frac{3}{2} \left| T \right| +\frac{1}{2}, -\frac{1}{2} \right\rangle &= \phi_8 \\
\left\langle -\frac{3}{2}, +\frac{3}{2} \left| T \right| +\frac{1}{2}, -\frac{1}{2} \right\rangle &= \phi_9,
\end{aligned}
\tag{A3}
$$

which are consistent with the definitions in Eq. (A2).

## Appendix B: Helicity Amplitudes of Photodisintegration in the Quark-Interchange Model

### *a. Quark Interchange model*

Following the approach presented for example in Refs.[42–45], the scattering amplitude for a process $ab \to cd$, in which $a, b, c$ and $d$ are baryons, is obtained from,

$$
\begin{aligned}
\langle cd \mid T \mid ab \rangle = & \sum_{\alpha,\beta,\gamma} \langle \psi_c^\dagger \mid \alpha_2', \beta_1', \gamma_1' \rangle \langle \psi_d^\dagger \mid \alpha_1', \beta_2', \gamma_2' \rangle \\
& \times \langle \alpha_2', \beta_2', \gamma_2', \alpha_1' \beta_1' \gamma_1' \mid H \mid \alpha_1, \beta_1, \gamma_1, \alpha_2 \beta_2 \gamma_2 \rangle \cdot \langle \alpha_1, \beta_1, \gamma_1 \mid \psi_a \rangle \langle \alpha_2, \beta_2, \gamma_2 \mid \psi_b \rangle,
\end{aligned}
\tag{B1}
$$

where $(\alpha_i, \alpha_i')$, $(\beta_i, \beta_i')$ and $(\gamma_i \gamma_i')$ describe the spin-flavor quark states before and after the hard scattering, $H$, and

$$
C^j_{\alpha,\beta,\gamma} = \langle \alpha, \beta, \gamma \mid \psi_j \rangle \tag{B2}
$$

describes the probability amplitude of finding an $\alpha, \beta, \gamma$ helicity-flavor combination of three valence quarks in the baryon $j$. These coefficients are obtained from the expansion of the

baryon's spin-isospin wave function in three-quark valence states as follows:

$$\begin{aligned}
\psi^{i^3_N,h_N} &= \frac{N}{\sqrt{2}}\Big\{\sigma(\chi^{(23)}_{0,0}\chi^{(1)}_{\frac{1}{2},h_N})\cdot(\tau^{(23)}_{0,0}\tau^{(1)}_{\frac{1}{2},i^3_N}) + \\
&\rho\sum_{i^3_{23}=-1}^{1}\sum_{h^3_{23}=-1}^{1}\langle 1,h_{23};\frac{1}{2},h_N-h_{23}\mid\frac{1}{2},h_N\rangle\langle 1,i^3_{23};\frac{1}{2},i^3_N-i^3_{23}\mid\frac{1}{2},i^3_N\rangle \\
&\times(\chi^{(23)}_{1,h_{23}}\chi^{(1)}_{\frac{1}{2},h_N-h_{23}})\cdot(\tau^{(23)}_{1,i^3_{23}}\tau^{(1)}_{\frac{1}{2},i^3_N-i^3_{23}})\Big\}. \qquad \text{(B3)}
\end{aligned}$$

The indexes 1 and 23 label the quark and the diquark states. The first term corresponds to quarks 2 and 3 being in a helicity zero isosinglet state, while the second term corresponds to quarks 2 and 3 in helicity 1-isotriplet states. Where $\chi$ and $\tau$ represent helicity and isospin states with helicity $h$ and isospin projection $i^3$ respectively. For the wave functions of $\Delta$-isobars $\sigma = 0$ and $\rho = 1$, while for nucleon wave functions $\sigma = 1$ and the parameter $\rho$ characterizes the average strength of the isotriplet diquark radial state relative to that of the isosinglet state. Two extreme values of $\rho = 1$ and $\rho = 0$ correspond to the realization of the SU(6) and good diquark symmetries in the wave function.

Using Eq.(B3) in Eq.(B1) for the hadronic scattering amplitude one obtains:

$$\langle cd|T^{QIM}|ab\rangle = A_{\alpha_1',\alpha_2',\alpha_1\alpha_2}(\theta^N_{c.m.})M^{ac}_{\alpha_1,\alpha_1'}M^{bd}_{\alpha_2,\alpha_2'} + A_{\alpha_1',\alpha_2',\alpha_1\alpha_2}(\pi-\theta^N_{c.m.})M^{ad}_{\alpha_1,\alpha_1'}M^{bc}_{\alpha_2,\alpha_2'}, \qquad \text{(B4)}$$

where

$$M^{ij}_{\alpha,\alpha'} = C^i_{\alpha,\beta,\gamma}C^j_{\alpha',\beta,\gamma} + C^i_{\beta,\alpha,\gamma}C^j_{\beta,\alpha',\gamma} + C^i_{\beta,\gamma,\alpha}C^j_{\beta,\gamma,\alpha'}, \qquad \text{(B5)}$$

which accounts for all possible interchanges of $\alpha$ and $\alpha'$ quarks leaving $\beta$ and $\gamma$ quarks unchanged. In the QI model the interchanging quarks conserve their corresponding helicities and flavors, this is accounted for in the matrix elements of $A$ in Eq.(B4.),

$$A_{\alpha_1',\alpha_2',\alpha_1\alpha_2}(s,\theta^N_{c.m.}) \propto \delta_{\alpha_1',\alpha_2}\delta_{\alpha_2',\alpha_1}\frac{f(\theta^N_{c.m.})}{s^2} \qquad \text{(B6)}$$

Eq.(B4) has two terms, first (referred as a $t$ term) in which four quarks scatter at angle $\theta^N_{c.m.}$ and two (interchanging) quarks scatter at $\pi - \theta^N_{c.m.}$ and the second (referred as a $u$ term) in which two interchanging quarks scatter at $\theta^N_{c.m.}$, while four spectator quarks scatter at $\pi - \theta^N_{c.m.}$.

### 1. Helicity Amplitudes in the Quark Interchange Model

Through the above procedure using Eq.(B4) for the helicity amplitudes of $pn$ scattering one obtains:

$$\begin{aligned}
\phi_1(\theta^N_{c.m.}) &= (2-y)f(\theta^N_{c.m.}) + (1+2y)f(\pi-\theta^N_{c.m.}) \qquad \text{(B7)}\\
\phi_2(\theta^N_{c.m.}) &= 0\\
\phi_3(\theta^N_{c.m.}) &= (2+y)f(\theta^N_{c.m.}) + (1+4y)f(\pi-\theta^N_{c.m.})\\
\phi_4(\theta^N_{c.m.}) &= 2yf(\theta^N_{c.m.}) + 2yf(\pi-\theta^N_{c.m.})\\
\phi_5(\theta^N_{c.m.}) &= 0,
\end{aligned}$$

were,

$$y = \frac{2}{3}\frac{\rho}{1+\rho^2}\left(1+\frac{2}{3}\frac{\rho}{1+\rho^2}\right). \tag{B8}$$

For $pn \to \Delta^+\Delta^0$ scattering amplitudes we obtain:

$$\begin{aligned}
\phi_1 &= \frac{2}{9}N_{\Delta\Delta}(2f(\theta^N_{c.m.}) - f(\pi-\theta^N_{c.m.})) \\
\phi_3 &= \frac{1}{9}N_{\Delta\Delta}(4f(\theta^N_{c.m.}) + f(\pi-\theta^N_{c.m.})) \\
\phi_4 &= \frac{2}{9}N_{\Delta\Delta}(f(\theta^N_{c.m.})) + f(\pi-\theta^N_{c.m.}) \\
\phi_6 &= \frac{N_{\Delta\Delta}}{3\sqrt{3}}(2f(\theta^N_{c.m.}) - f(\pi-\theta^N_{c.m.})) \\
\phi_7 &= \frac{N_{\Delta\Delta}}{3\sqrt{3}}(2f(\theta^N_{c.m.}) - f(\pi-\theta^N_{c.m.})) \\
\phi_8 &= \frac{2}{9}N_{\Delta\Delta}f(\theta^N_{c.m.}) \\
\phi_9 &= \frac{1}{3}N_{\Delta\Delta}f(\pi-\theta^N_{c.m.}),
\end{aligned} \tag{B9}$$

and similarly for the amplitudes of the $pn \to \Delta^{++}\Delta^-$ scattering, QI model gives:

$$\begin{aligned}
\phi_1 &= -\frac{2}{3}N_{\Delta\Delta}f(\theta^N_{c.m.}) \\
\phi_3 &= -\frac{2}{3}N_{\Delta\Delta}f(\theta^N_{c.m.}) \\
\phi_4 &= -\frac{1}{3}N_{\Delta\Delta}f(\theta^N_{c.m.}) \\
\phi_6 &= \frac{-N_{\Delta\Delta}}{\sqrt{3}}f(\theta^N_{c.m.}) \\
\phi_7 &= \frac{-N_{\Delta\Delta}}{\sqrt{3}}f(\theta^N_{c.m.}) \\
\phi_8 &= -N_{\Delta\Delta}f(\theta^N_{c.m.}) \\
\phi_9 &= 0,
\end{aligned} \tag{B10}$$

For both sets of equations in (B9) and (B10), we have

$$N_{\Delta\Delta} = \frac{(1+\rho)^2}{1+\rho^2}, \tag{B11}$$

which shows that the strength of the two $\Delta\Delta$ channels relative to each other is independent of the value of $\rho$. This is not the case for their strengths relative to the $pn$ channel; from Eqs. (B8) we see that the $\rho$ dependence of the helicity amplitudes in $pn \to pn$ cannot be factorized.

### 2. Quark-Charge Factors

In the hard rescattering model, photodisintegration amplitudes are expressed in terms of hadronic scattering amplitudes weighted by the charges of struck quarks, Eq.(5). We further split the amplitude of Eq.(5) into $t$ and $u$ channel scatterings:

$$\sum_i Q_i^{N_k} \langle \lambda_{2f}; \lambda_{1f} \mid T_{(pn \to B_1 B_2),i}(s, \tilde{t}) \mid \lambda_\gamma; \lambda_{2i} \rangle = \left[ Q_j^{tN_k} \phi_j^t + Q_j^{uN_k} \phi_j^u \right], \tag{B12}$$

where $Q_i^{t/uN}$ is the charge of the quark, struck by the incoming photon from the nucleon $N$ with further $\theta^N_{c.m.}$ or $\pi - \theta^N_{c.m.}$ scattering. The helicity amplitudes are also split into $t$ and $u$ parts

$$\begin{aligned} \phi_i(\theta^N_{c.m.}) &= \phi_i^t(\theta^N_{c.m.}) + \phi_i^u(\theta^N_{c.m.}) \\ &= c_t f(\theta^N_{c.m.}) + c_u f(\pi - \theta^N_{c.m.}), \end{aligned} \tag{B13}$$

with $\phi^t$ and $\phi^u$ corresponding to the $\theta^N_{c.m.}$ or $\pi - \theta^N_{c.m.}$ scattering terms in Eq.(B4).

Using the above definitions and Eqs.(B4,B5,B8,B9,B10) the charge factors $Q^t$ and $Q^u$ are calculated using the following relations:

$$\begin{aligned} Q_j^{tN_k} &= \frac{Q(\alpha_k) A_{\alpha_1',\alpha_2',\alpha_1\alpha_2} M^{ac}_{\alpha_1,\alpha_1'} M^{bd}_{\alpha_2,\alpha_2'}}{\phi_j^t} \\ Q_j^{uN_k} &= \frac{Q(\alpha_k) A_{\alpha_1',\alpha_2',\alpha_1\alpha_2} M^{ad}_{\alpha_1,\alpha_1'} M^{bc}_{\alpha_2,\alpha_2'}}{\phi_j^u}, \end{aligned} \tag{B14}$$

where summation is understood for repeated $\alpha$ indices, $Q(\alpha)$ is the charge in e units of a quark $\alpha$ and the index $j$ labels the process $ab \to cd$.


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